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Lowering α_s by flipping SU(5)

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Abstract

We show that the prediction for $\alpha_s(M_Z)$ in flipped SU(5) is naturally lower than in minimal SU(5), and that the former can accommodate the full range of $\alpha_s(M_Z)$ presently allowed by experiment. Our computations include two-loop ($\delta_{2\text{loop}}$) and light (δ_{light}) and heavy (δ_{heavy}) threshold effects. Unlike minimal SU(5), in flipped SU(5) the heavy threshold effects can naturally decrease the predicted value of $\alpha_s(M_Z)$. We also show that the value of the proton lifetime into the dominant channel $p \rightarrow e^+ \pi^0$ is within the observable range at SuperKamiokande, and should discriminate against minimal supersymmetric SU(5), where the dominant mode is $p \rightarrow \bar{\nu} K^+$.

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One of the most impressive pieces of circumstantial evidence for Grand Unified Theories is the successful correlation they yield between $\alpha_s(M_Z)$ and $\sin^2 \theta_W$ [1]. Historically, this was first phrased as a prediction of $\sin^2 \theta_W$ based on the measured value of α_s , whose qualitative agreement with early electroweak data was impressive, particularly for supersymmetric GUTs [2]. The advent of precision electroweak data from LEP and elsewhere [3], and the measurement of the top-quark mass m_t [4] has enabled $\sin^2 \theta_W$ to be determined with such accuracy that it is natural to turn the GUT correlation round the other way, and use it to predict $\alpha_s(M_Z)$, which is still not known with satisfactory precision [5]. The prediction of a minimal non-supersymmetric SU(5) GUT is disastrously low, and debate centres on the the possibility of discriminating between different supersymmetric GUTs and/or obtaining indirect constraints on the scale of supersymmetry breaking.

Progress on these questions is hampered by the persistent imprecision in the experimental value of $\alpha_s(M_Z)$, which is usually quoted as 0.118 ± 0.006 . Most determinations lie within one standard deviation of this mean, but this consistency hides two schools of thought: one supported more strongly by high-energy determinations of α_s which favour $\alpha_s(M_Z) \geq 0.120$ [3], and one favoured more strongly by low-energy determinations of α_s which favour $\alpha_s(M_Z) \leq 0.115$ when evolved up to a higher scale [6]. It is not clear whether this apparent discrepancy is real, but it has led to various supersymmetric speculations, including the possibility that gluinos are sufficiently light to alter the evolution of α_s between low energies and M_Z [7], or that radiative corrections due to other light particles such as the chargino and top-squark reduce the value of $\alpha_s(M_Z)$ inferred from LEP data [8].

The minimal supersymmetric SU(5) prediction of $\alpha_s(M_Z)$ definitely lies on the upper side of the experimental range, and even above it, with values around 0.130 or higher being preferred [9, 10]. This preference should however be treated with caution, as the predicted value may be reduced by GUT threshold effects [11, 12] or by ‘slop’ induced by Planck-scale interactions [13]. One example of a supersymmetric GUT which makes a lower prediction for $\alpha_s(M_Z)$ is the Missing-Doublet Model (MDM) [14], in which the GUT threshold corrections are generally negative [15]. However, the MDM is quite cumbersome, containing several large Higgs representations that cannot easily be accommodated in a string framework [16].

By far the most economical realization of the missing-partner mechanism is that in flipped SU(5) [17], which needs only **10** and $\overline{\mathbf{10}}$ representations of GUT Higgses and has been derived from string [18]. This model also provides another natural mechanism for reducing the prediction for $\alpha_s(M_Z)$, namely if the scale at which α_s and the SU(2) electroweak coupling become equal is lower than the scale at which the SU(5) and U(1) couplings become equal [19].

In this paper we make a quantitative exploration of this possibility in the supersymmetric minimal flipped SU(5) GUT [17]. In addition to exploring the above-

mentioned possible difference in unification scales, we also calculate for the first time the GUT threshold effects in flipped SU(5), finding that they may be naturally negative. We also demonstrate that the flipped SU(5) mechanism for reducing $\alpha_s(M_Z)$ has an observable signature if it lies in the lower part of the range currently allowed by experiment [5], namely that proton decay into $e^+\pi^0$ should be observable in the SuperKamiokande detector.

The starting point for our discussion is the lowest-order prediction for $\alpha_s(M_Z)$ in SU(5) GUTs, namely

$$\alpha_s(M_Z) = \frac{\frac{7}{3}\alpha}{5\sin^2\theta_W - 1} , \quad (1)$$

which yields the present central experimental value of $\alpha_s(M_Z) \approx 0.118$ for $\sin^2\theta_W = 0.231$ and $\alpha^{-1} = 128$. It is possible to include two-loop corrections ($\delta_{2\text{loop}}$), and light (δ_{light}) and heavy (δ_{heavy}) threshold corrections by making the following substitution in Eq. (1) [19]:

$$\sin^2\theta_W \rightarrow \sin^2\theta_W - \delta_{2\text{loop}} - \delta_{\text{light}} - \delta_{\text{heavy}} . \quad (2)$$

Here $\delta_{2\text{loop}} \approx 0.0030$ whereas δ_{light} and δ_{heavy} can have either sign. Neglecting δ_{light} and δ_{heavy} , the SU(5) prediction increases to $\alpha_s(M_Z) \approx 0.130$. Therefore, if one wishes to obtain a value of $\alpha_s(M_Z)$ within one standard deviation of the present central experimental value, one requires a non-negligible contribution from δ_{light} and/or δ_{heavy} , so that the combined correction ($\delta_{2\text{loop}} + \delta_{\text{light}} + \delta_{\text{heavy}}$) is suppressed. In large regions of parameter space, $\delta_{\text{light}} > 0$ and does not help. As we discuss in more detail later, exploiting δ_{heavy} is difficult in minimal SU(5) because of proton decay constraints [11, 10], unless one goes to the limit $m_0 \gg m_{1/2}$ and $m_0 \sim 1$ TeV which suppresses proton decay. However, even in this case the decrease in $\alpha_s(M_Z)$ is not large: $\alpha_s(M_Z) > 0.123$ [10]. One way to circumvent these problems is the cumbersome missing doublet model (MDM) [10, 15]: here we take the more elegant route of flipped SU(5).

In this model, there is a first unification scale M_{32} at which the SU(3) and SU(2) gauge couplings become equal, given to lowest order by [20]

$$\frac{1}{\alpha_2} - \frac{1}{\alpha_5} = \frac{b_2}{2\pi} \ln \frac{M_{32}}{M_Z} , \quad (3)$$

$$\frac{1}{\alpha_3} - \frac{1}{\alpha_5} = \frac{b_3}{2\pi} \ln \frac{M_{32}}{M_Z} , \quad (4)$$

where $\alpha_2 = \alpha/\sin^2\theta_W$, $\alpha_3 = \alpha_s(M_Z)$, and the one-loop beta functions are $b_2 = +1$, $b_3 = -3$. On the other hand, the hypercharge gauge coupling $\alpha_Y = \frac{5}{3}(\alpha/\cos^2\theta_W)$ evolves in general to a different value α'_1 at the scale M_{32} :

$$\frac{1}{\alpha_Y} - \frac{1}{\alpha'_1} = \frac{b_Y}{2\pi} \ln \frac{M_{32}}{M_Z} , \quad (5)$$

with $b_Y = \frac{33}{5}$. Above the scale M_{32} the gauge group is $SU(5) \times U(1)$, with the $U(1)$ gauge coupling α_1 related to α'_1 and the $SU(5)$ gauge coupling (α_5) by

$$\frac{25}{\alpha'_1} = \frac{1}{\alpha_5} + \frac{24}{\alpha_1} . \quad (6)$$

The $SU(5)$ and $U(1)$ gauge couplings continue to evolve above the scale M_{32} , eventually becoming equal at a higher scale M_{51} . The consistency condition that $M_{51} \geq M_{32}$ requires $\alpha_1(M_{32}) \leq \alpha_5(M_{32})$ which, according to Eq. (6), also implies $\alpha'_1 \leq \alpha_5(M_{32})$. The maximum possible value of M_{32} , namely M_{32}^{\max} , is obtained when $\alpha'_1 = \alpha_5(M_{32})$ and is given by the following relation obtained from Eq. (5):

$$\frac{1}{\alpha_Y} - \frac{1}{\alpha_5} = \frac{b_Y}{2\pi} \ln \frac{M_{32}^{\max}}{M_Z} , \quad (7)$$

which coincides with the unification scale in minimal $SU(5)$. We thus obtain a set of three equations (3,4,7) involving six variables: $\alpha_2, \alpha_3, \alpha_Y$ and $\alpha_5, M_{32}^{\max}, M_{32}$. Solving these equations for the value of $\alpha_s(M_Z)$, the analogue of Eq. (1) is given by

$$\alpha_s(M_Z) = \frac{\frac{7}{3}\alpha}{5 \sin^2 \theta_W - 1 + \frac{11}{2\pi} \alpha \ln(M_{32}^{\max}/M_{32})} . \quad (8)$$

It is clear that the prediction for $\alpha_s(M_Z)$ in flipped $SU(5)$ is automatically smaller than in minimal $SU(5)$.

The next-to-leading order corrections to Eq. (8) are also obtained by the substitution in Eq. (2). Numerically, an increase of $\sim 10\%$ in the denominator in Eq. (1), which would compensate for the decrease due to $\delta_{2\text{loop}}$, is achieved simply by setting $M_{32} \approx \frac{1}{3} M_{32}^{\max}$ in Eq. (8). We also note that the value of α_5 at M_{32} is given by

$$\frac{1}{\alpha_5} = \frac{1}{\alpha_5^{\max}} + \frac{33}{28} \frac{1}{2\pi} \ln \frac{M_{32}^{\max}}{M_{32}} , \quad (9)$$

where α_5^{\max} is the maximum possible value of α_5 , which is that attained in the minimal $SU(5)$ case.

The above qualitative discussion shows the differences between flipped $SU(5)$ and minimal $SU(5)$. We now perform a detailed numerical calculation of $\delta_{2\text{loop}}$, δ_{light} and M_{32}^{\max} , using the approach of Refs. [11, 21]. The δ_{light} correction is given in the step-function approximation by

$$\begin{aligned} \delta_{\text{light}} = & \frac{\alpha}{20\pi} \left[-3L(m_t) + \frac{28}{3}L(m_{\tilde{g}}) - \frac{32}{3}L(m_{\tilde{w}}) - L(m_h) - 4L(m_H) \right. \\ & \left. + \frac{5}{2}L(m_{\tilde{q}}) - 3L(m_{\tilde{\ell}_L}) + 2L(m_{\tilde{\ell}_R}) - \frac{35}{36}L(m_{\tilde{t}_1}) - \frac{19}{36}L(m_{\tilde{t}_2}) \right] , \end{aligned} \quad (10)$$

where $L(x) = \ln(x/M_Z)$. In the limit where the $L(x)$ are large, the sparticle spectrum may be parametrized approximately in terms of the universal soft supersymmetry-breaking scalar (m_0) and gaugino ($m_{1/2}$) masses: $m_{\tilde{g}} = 2.7m_{1/2}$, $m_{\tilde{w}} = 0.79m_{1/2}$,

$m_i^2 = m_0^2 + c_i^2 m_{1/2}^2$, with $c_{\tilde{q}} = 6$, $c_{\tilde{t}_L} = 0.5$, $c_{\tilde{t}_R} = 0.15$. In the large- L limit we may also consider the Higgs spectrum as containing a lighter doublet with mass M_Z , and a heavier Higgs doublet whose mass is identified with the Higgs mixing parameter μ ($m_H = |\mu|$). Finally, the top-squark masses are parametrized as $m_{\tilde{t}_{1,2}}^2 = m_{\tilde{q}}^2 \pm \bar{m}^2$, where $\bar{m} \leq m_{\tilde{q}}$ accounts for the splitting of the top-squark mass eigenstates, and is in principle related to the trilinear scalar coupling A . We scan the four-dimensional parameter space $(m_0, m_{1/2}, \mu, \bar{m})$ in the region below 1 TeV, but do not attempt to incorporate either the radiative electroweak breaking constraint or any fine-tuning condition, which would in any case lead to correlations among the mass parameters, and thus to a more restrictive range of possible values of δ_{light} .

The effect (δ_{heavy}) of the GUT thresholds is well known in minimal SU(5), where it yields the expression [11]

$$\delta_{\text{heavy}}^{\text{SU}(5)} = \frac{\alpha}{20\pi} \left[-6 \ln \frac{M_U}{M_{H_3}} + 4 \ln \frac{M_U}{M_V} + 2 \ln \frac{M_U}{M_\Sigma} \right], \quad (11)$$

where $M_U = \max\{M_{H_3}, M_V, M_\Sigma\}$. In this expression, M_V represents the mass of the X, Y gauge bosons and gauginos, and the Higgs bosons and Higgsinos in the **24**; M_Σ represents the mass of the remaining states in the **24** (a color octet and an SU(2) triplet); and M_{H_3} represents the common mass of the triplet components of the Higgs pentaplets. Since $M_{H_3} > M_V$ is required to suppress dimension-five proton decay operators in minimal SU(5) [22, 23], $\delta_{\text{heavy}}^{\text{SU}(5)} > 0$ in general [11], unless one resorts to unnaturally large sparticle masses.

The corresponding expression for δ_{heavy} in flipped SU(5) can be obtained from the general formula [11]

$$\delta_{\text{heavy}} = -\frac{\alpha}{20\pi} \sum_{R_i} C(R_i) \ln \frac{M_U}{M_i}, \quad (12)$$

where the sum runs over all GUT scale particles R_i with masses M_i , and

$$C(R) = \frac{10}{3}b_Y(R) - 8b_2(R) + \frac{14}{3}b_3(R), \quad (13)$$

is a linear combination of the one-loop contributions to the beta functions. The spectrum of heavy particles in flipped SU(5) consists of the X, Y gauge supermultiplets, the **10** and $\overline{\mathbf{10}}$ Higgs supermultiplets $H = \{Q_H, d_H^c, \nu_H^c\}$ and $\bar{H} = \{Q_{\bar{H}}, d_{\bar{H}}^c, \nu_{\bar{H}}^c\}$ that break SU(5), and the Higgs triplets (h_3, \bar{h}_3) contained in the Higgs pentaplets (h, \bar{h}) .¹ In the SU(5) symmetry breaking process, the neutral components of the **10** and $\overline{\mathbf{10}}$ Higgs multiplets acquire equal vevs $\langle \nu_H^c \rangle = \langle \nu_{\bar{H}}^c \rangle = V$, giving equal masses ($M_V = g_5|V|$) to the X, Y gauge bosons and gauginos, and to the $Q_H, Q_{\bar{H}}$ Higgs boson and Higgsinos. The remaining components ($d_H^c, d_{\bar{H}}^c$) of the **10** and $\overline{\mathbf{10}}$ participate

¹There are also several gauge-singlet representations, such as the right-handed neutrinos, which have large masses but do not contribute to δ_{heavy} , since they have $C(R) = 0$.

in the economical missing partner mechanism of flipped SU(5), combining with the h_3, \bar{h}_3 triplets in the Higgs pentaplets (h, \bar{h}) via the superpotential couplings $\lambda_4 H H h$ and $\lambda_5 \bar{H} \bar{H} \bar{h}$ to acquire masses $M_{H_3} = \lambda_4 |V|$ and $M_{\bar{H}_3} = \lambda_5 |V|$. We note that, unlike the case of minimal SU(5), there are no contributions from Σ -type fields since these are absent from the spectrum of the model. (We recall that the ν_H^c and $\nu_{\bar{H}}^c$ are gauge singlets, and hence do not contribute to δ_{heavy}). Calculating the $C(R)$ coefficients in Eq. (12), we obtain

$$\delta_{\text{heavy}} = \frac{\alpha}{20\pi} \left[-6 \ln \frac{M_{32}}{M_{H_3}} - 6 \ln \frac{M_{32}}{M_{\bar{H}_3}} + 4 \ln \frac{M_{32}}{M_V} \right] = \frac{\alpha}{20\pi} \left[-6 \ln \frac{r^{4/3} g_5^{2/3}}{\lambda_4 \lambda_5} \right] \quad (14)$$

where $r = \max\{g_5, \lambda_4, \lambda_5\}$. Since the H_3 and \bar{H}_3 do not mix, they do not contribute significantly to proton decay, and hence there is no strong constraint on M_{H_3, \bar{H}_3} from proton decay in flipped SU(5). We therefore see from Eq. (14) that it is possible that $M_{H_3, \bar{H}_3} < M_V = M_U$ (*i.e.*, $r = g_5$), in which case $\delta_{\text{heavy}} < 0$ naturally. For instance, if $\lambda_4, \lambda_5 \sim \frac{1}{8}g_5$, then $\delta_{\text{heavy}} \approx -0.0030$, which completely compensates the δ_{2loop} contribution.

In Fig. 1 we show the prediction for $\alpha_s(M_Z)$ in flipped SU(5) as a function of the ratio $M_{32}/M_{32}^{\text{max}}$. The solid lines delimit the range of predictions obtained using the latest experimental value $\sin^2 \theta_W = 0.23143 \pm 0.00028$ and setting $\delta_{\text{light}} + \delta_{\text{heavy}} = 0$. We note that the minimal SU(5) prediction is obtained when $M_{32}/M_{32}^{\text{max}} = 1$. Scanning over the possible range of sparticle masses described above, we find significant variations in the predictions for $\alpha_s(M_Z)$, indicated by the dashed lines in Fig. 1, usually towards higher values (*i.e.*, $\delta_{\text{light}} > 0$). The maximum values of $\alpha_s(M_Z)$ are obtained for $m_0 \gg m_{1/2}, |\mu|$, whereas the minimum values are obtained for $|\mu| \gg m_0, m_{1/2}$.² The general detrimental effect of δ_{light} may be compensated somewhat by $\delta_{\text{heavy}} < 0$ effects.

In the case of minimal SU(5) (*i.e.*, $M_{32}/M_{32}^{\text{max}} = 1$), these effects may decrease the value of $\alpha_s(M_Z)$ as low as 0.123 [10]. However, in flipped SU(5) one may find $\alpha_s(M_Z) \lesssim 0.120$ even without resorting to such stratagems, if the unification scale is lowered enough. Equations (2) and (8) show that including the effects of δ_{heavy} simply amounts to a re-scaling of the $M_{32}/M_{32}^{\text{max}}$ axis on Fig. 1, *i.e.*,

$$\frac{M_{32}}{M_{32}^{\text{max}}} \rightarrow \frac{M_{32}}{M_{32}^{\text{max}}} e^{-10\pi \delta_{\text{heavy}}/11\alpha} . \quad (15)$$

For definiteness, let us assume that $\frac{1}{\rho} \lesssim g_5/\sqrt{\lambda_4 \lambda_5} \lesssim \rho$, with the heuristic choice $\rho = 3$. Equation (14) then implies $0.0005 \gtrsim \delta_{\text{heavy}} \gtrsim -0.0016$, and the multiplicative factor in Eq. (15) is in the range $0.81 \lesssim e^{-10\pi \delta_{\text{heavy}}/11\alpha} \lesssim 1.8$.

²Note, however, that this scenario is disfavored by radiative symmetry-breaking constraints, which imply $|\mu| \sim \max\{m_0, m_{1/2}\}$.

We now turn to a possible experimental signature of this flipped SU(5) mechanism for lowering α_s , namely proton decay, which is clearly enhanced by decreasing the unification scale. As is well known, one of the attractive features of flipped SU(5) is the fact that dimension-five proton decay operators are highly suppressed in this model [20]. However, the usual dimension-six proton decay operators which would lead, for example, to $p \rightarrow e^+\pi^0$ decay, are present and their coefficients would be enhanced relative to those in minimal SU(5) by a factor of $(M_{32}^{\max}/M_{32})^4$. Concretely, we estimate that [20, 24]

$$\tau(p \rightarrow e^+\pi^0) \approx 1.5 \times 10^{33} \left(\frac{M_{32}}{10^{15} \text{ GeV}} \right)^4 \left(\frac{0.042}{\alpha_5} \right)^2, \quad (16)$$

where we have taken the central value of the $p \rightarrow e^+\pi^0$ matrix element from [25], which quotes an error of about 20% in the decay rate. The error quoted in [25] is mainly statistical in nature, and does not include possible systematic errors associated with the quenched-fermion approximation, finite lattice size and spacing, etc. We consider it prudent to double the quoted error when interpreting (16) in order to allow for such possible effects. This would, however, correspond to a variation in M_{32} of no more than 10%.

Equipped with equation (16), we plot the prediction for $\alpha_s(M_Z)$ versus $\tau(p \rightarrow e^+\pi^0)$ in Fig. 2, indicating the effect of δ_{light} . As discussed above, for a fixed value of $\alpha_s(M_Z)$, heavy threshold corrections introduce a multiplicative uncertainty in M_{32} in the range 0.8–1.8, and thus a multiplicative uncertainty in the proton lifetime in the range 0.4–10. It is interesting to note that the present experimental lower bound $\tau(p \rightarrow e^+\pi^0)^{\text{exp}} > 5.5 \times 10^{32} \text{ y}$ [26] allows values of $\alpha_s(M_Z)$ as low as 0.108, even without accounting for the uncertainties in the proton decay calculation. This approximate lower bound is indicated by the crosses in Fig. 1. Moreover, for values of $\alpha_s(M_Z) \lesssim 0.114$, the mode $p \rightarrow e^+\pi^0$ should be observable at SuperKamiokande in flipped SU(5), whereas it is expected to be utterly unobservable in minimal supersymmetric SU(5). Furthermore, even though flipped SU(5) and minimal supersymmetric SU(5) predict similar proton decay modes, the relative branching ratios differ, because of the different operator structures in the two models. We also recall that the mode $p \rightarrow \bar{\nu}K^+$ via dimension-five operators should be dominant in minimal supersymmetric SU(5) [22].

We conclude by putting the present investigation in the proper perspective regarding flipped SU(5) model building. As is well known, thanks to the simple “flip” $e^c \leftrightarrow \nu^c$, flipped SU(5) can be broken by vevs of small representations (**10**, $\overline{\mathbf{10}}$) in a unique way. The presence of ν^c leads to a see-saw mechanism for generating small neutrino masses, and possibly a cosmological baryon asymmetry which is later recycled into a baryon asymmetry by sphaleron interactions [27]. The triplet components of the **10**, $\overline{\mathbf{10}}$ yield a very economical doublet-triplet splitting, which in turn suppresses greatly dimension-five proton decay operators. This natural suppression

allows the heavy Higgs triplets to contribute negatively to δ_{heavy} , so that the prediction for $\alpha_s(M_Z)$ can be lowered to anywhere in the 0.11-0.12 interval. If in the future the low- and high-energy determinations of $\alpha_s(M_Z)$ are reconciled experimentally with a value in the lower half of the range presently allowed, the SuperKamiokande detector should observe $p \rightarrow e^+\pi^0$ decays, whereas minimal supersymmetric SU(5) would predict $p \rightarrow \bar{\nu}K^+$.

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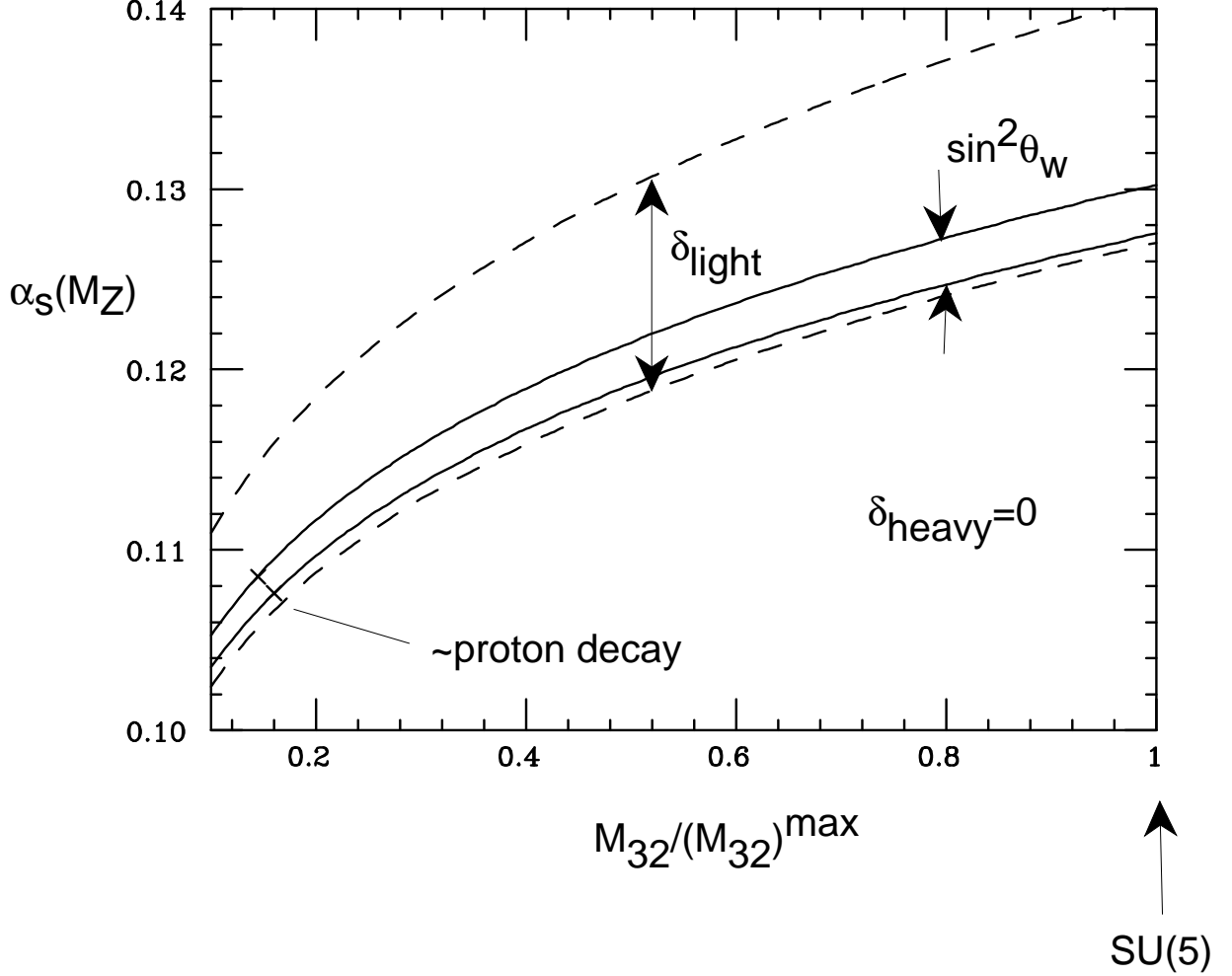


Figure 1: The prediction for $\alpha_s(M_Z)$ in flipped SU(5) as a function of the ratio M_{32}/M_{32}^{\max} . The minimal SU(5) predictions are recovered for $M_{32}/M_{32}^{\max} = 1$. The solid lines represent the range of predictions for $\sin^2 \theta_W = 0.23143 \pm 0.00028$ with no threshold corrections ($\delta_{\text{light}} = \delta_{\text{heavy}} = 0$). The dashed lines represent the excursion obtained by scanning over sparticle masses ($\delta_{\text{light}} \neq 0$) below 1 TeV. Inclusion of typical δ_{heavy} values amounts to a rescaling of the M_{32}/M_{32}^{\max} axis by a factor in the range 0.8–1.8. The crosses indicate approximate lower bounds from proton decay constraints, which are displayed in more detail in Figure 2.

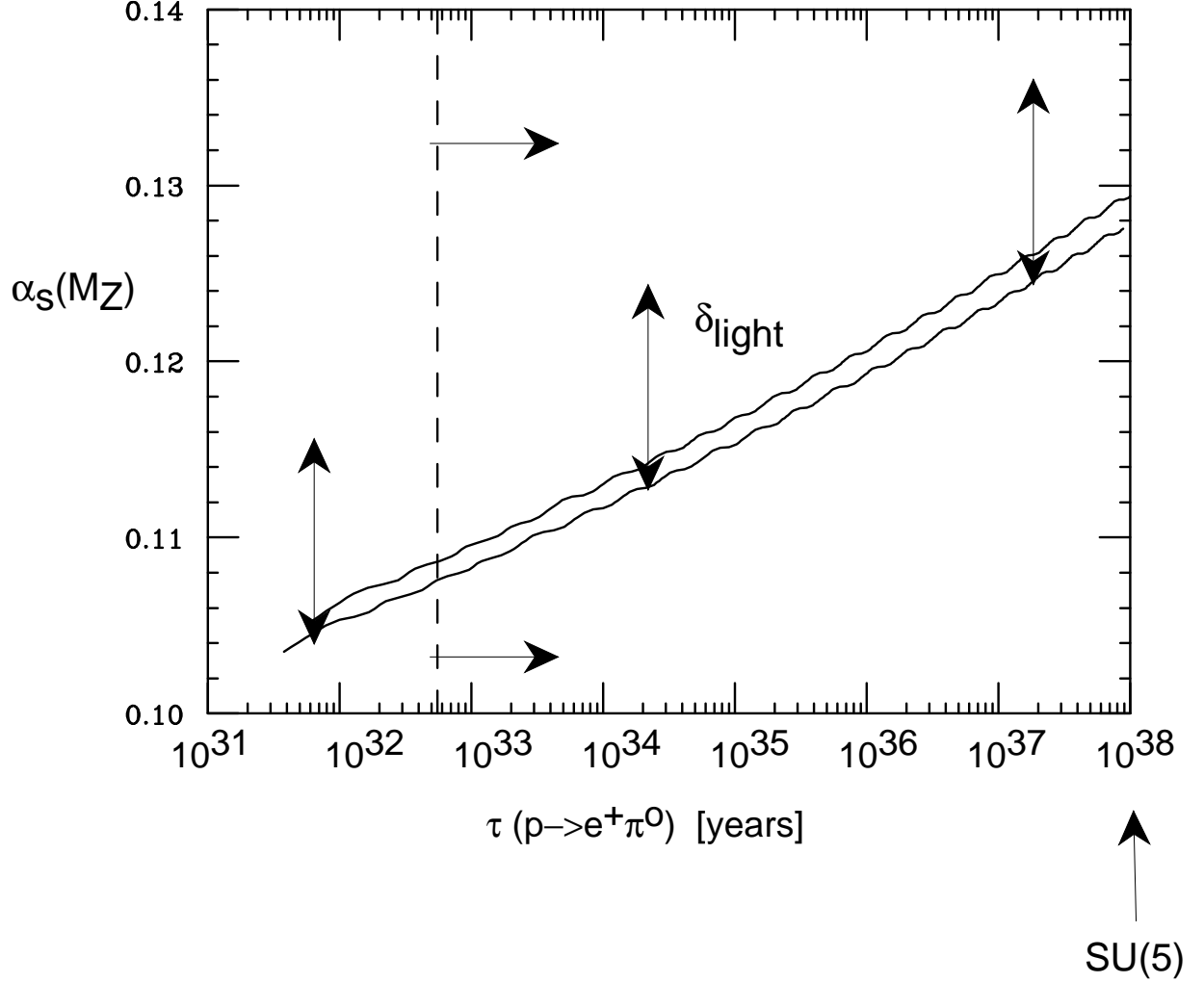


Figure 2: The prediction for $\alpha_s(M_Z)$ in flipped $\text{SU}(5)$ as a function of the proton lifetime into the mode $e^+ \pi^0$. The present experimental lower bound is indicated by the dashed line. The minimal $\text{SU}(5)$ predictions are recovered for $M_{32}/M_{32}^{\text{max}} = 1$. The solid lines represented the range of predictions for $\sin^2 \theta_W = 0.23143 \pm 0.00028$ with no threshold corrections ($\delta_{\text{light}} = \delta_{\text{heavy}} = 0$). The arrows indicate the extent of the excursion obtained by scanning over sparticle masses ($\delta_{\text{light}} \neq 0$) below 1 TeV. Heavy threshold corrections amount to a multiplicative uncertainty factor in the proton lifetime in the range 0.4–10.